

SEARCHING FOR SUSY DARK MATTER

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ABSTRACT

The possibility of detecting supersymmetric dark matter is examined within the framework of the minimal supergravity model (MSGM), where the \tilde{Z}_1 is the LSP for almost the entire parameter space. A brief discussion is given of experimental strategies for detecting dark matter. The relic density is constrained to obey $0.10 \leq \Omega_{\tilde{Z}_1} h^2 \leq 0.35$, consistent with COBE data. Expected event rates for an array of possible terrestrial detectors (^3He , CaF_2 , Ge , GaAs , NaI and Pb) are examined. In general, detectors relying on coherent \tilde{Z}_1 -nucleus scattering are more sensitive than detectors relying on incoherent (spin-dependent) scattering. The dependence of the event rates as a function of the SUSY parameters are described. The detectors are generally most sensitive to the small m_0 and small $m_{\tilde{q}}$ and large $\tan\beta$ part of the parameter space. The current $b \rightarrow s + \gamma$ decay rate eliminates regions of large event rates for $\mu > 0$, but allows large event rates to still occur for $\mu < 0$. MSGM models that also possess SU(5)-type proton decay generally predict event rates below the expected sensitivity of current dark matter detectors.

1. Introduction

If the SUSY models currently being examined are correct, then supersymmetry will be discovered at the LHC in the year 2005 or possibly even at an upgraded version of the Tevatron (e.g. the DiTevatron) in the year 2000. However, high energy colliders may not shed further light until then. Thus it is of interest to look at other phenomena which supersymmetry might effect, e.g. dark matter, proton decay, the $b \rightarrow s + \gamma$ decay etc. Each of these restricts the parameter space of supersymmetry, and so by combining the constraints, one can get sharper predictions of what to expect at colliders. We will discuss here the question of detecting SUSY dark matter, and how bounds on other processes affect dark matter searches. There is a warning however one should make concerning such analyses: In applying SUSY to dark matter searches, one is making *additional* assumptions

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(e.g., cosmological assumptions) not made in dealing with accelerator phenomena. We will test the sensitivity of final results to some of these extra assumptions, but some caution is needed in interpreting the theoretical predictions.

2. Dark Matter

There is much astronomical evidence that more than 90% of the total mass of the Galaxy and perhaps of the universe is made up of dark matter of an unknown type. In the Galaxy this can be seen from rotation curves of luminous matter (Fig. 1). The circular velocity v_{cir} does not fall with r beyond the optical radius. Similar

Fig. 1. Estimated mean spherical density of dark matter in the Galaxy¹.

effects are seen in other galaxies (Fig. 2). In the vicinity of the sun, the mean density of dark matter (DM) is estimated as

$$\rho_{DM} \cong 0.3 \text{ GeV}/cm^3 \quad (1)$$

(about $10^4 \gamma_{universe}$) and assuming a Maxwell velocity distribution, the DM has a velocity relative to the solar system of

$$v_{DM} \cong 320 km/s \quad (2)$$

i.e. $v_{DM}/c \approx 10^{-3}$.

Fig. 2 v_{cir} for a number of galaxies showing v_{cir} remains approximately constant well beyond the optical radius².

There are a large number of candidates for dark matter, both from astronomy and particle physics. Supersymmetric models with R-parity offer two candidates: the lightest neutralino \tilde{Z}_1 and the sneutrino, $\tilde{\nu}$. However, in supergravity models almost always the \tilde{Z}_1 is the lightest supersymmetric particle (LSP) and hence is absolutely stable. Thus the relic \tilde{Z}_1 , left over from the big bang would be the dark matter. The dynamics is then well fixed and we will deal with this case exclusively here.

COBE data suggests that DM is a mix of cold dark matter (CDM) [which we are assuming here to be the \tilde{Z}_1] and hot dark matter (HDM) [possibly massive neutrinos] in the ratio of about 2:1. There can also be baryonic dark matter (B) but nucleosynthesis analyses limit this to $\lesssim 10\%$. Thus if we define $\Omega_i = \rho_i/\rho_c$, where ρ_i is the mass density of type i and $\rho_c = 3H^2/(8\pi G_N)$ is the critical mass density to close the universe (H =Hubble constant and G_N =Newtonian gravitational constant), the inflationary scenario requires $\Sigma\Omega_i = 1$ and hence a reasonable mix of matter is

$$\Omega_{\tilde{Z}_1} \simeq 0.6; \quad \Omega_{HDM} \simeq 0.3; \quad \Omega_B \simeq 0.1 \quad (3)$$

What can be calculated theoretically is $\Omega_{\tilde{Z}_1} h^2$ where $h=H/(100 \text{ km/s Mpc})$. Astronomical observations give a range of values for h i.e. $h \cong 0.5 - 0.75$. Hence

$$\Omega_{\tilde{Z}_1} h^2 \cong 0.10 - 0.35 \quad (4)$$

This bound strongly restricts the SUSY parameter space. We will discuss below how sensitive our results are to the endpoints of this bound on $\Omega_{\tilde{Z}_1} h^2$.

3. Detection Strategies

The solar system is presumably being bombarded with \tilde{Z}_1 particles moving with velocity $< v_{\tilde{Z}_1} > \cong 320$ km/s. Two strategies have been proposed for their detection.

3.1 Indirect Detection

The \tilde{Z}_1 impinging on the sun can be relatively easily captured since the escape velocity at the surface of the sun is 618 km/s $\cong v_{\tilde{Z}_1}$. Once captured the \tilde{Z}_1 will be slowed down inside the sun and gradually fall to the center. There they accumulate and can annihilate giving rise to neutrinos via e.g. $\tilde{Z}_1 + \tilde{Z}_1 \rightarrow b + \bar{b}, \dots \rightarrow \nu_\mu + X$. Since $m_{\tilde{Z}_1} \cong (10-100)$ GeV, high energy neutrinos coming from the sun from this process would be a striking signal that could be observed on Earth by a neutrino telescope. Calculations indicate that one would need a telescope of area $> 1 km^2$ to cover the SUSY parameter space³, and telescopes of this size are currently being built.

3.2 Direct Detection

A direct approach to see the incoming \tilde{Z}_1 is to detect their scattering by quarks in nuclei of a terrestrial detector: $\tilde{Z}_1 + q \rightarrow \tilde{Z}_1 + q$. Two types of detectors being considered are the following.

3.21 Low Temperature Detectors

Since $m_{\tilde{Z}_1} \simeq (10-100)$ GeV, the recoil energy to a nucleus that has been struck is $\Delta E \approx (v_{\tilde{Z}_1}/c)^2 (m_{\tilde{Z}_1} c^2) \simeq (10-100)$ KeV. This is of a size to produce phonons (heat) or ionization in the lattice of the detector. The temperature rise ΔT is $\Delta T = \Delta E / VC$ where C is the specific heat and V the volume. Since at low temperature $C \sim T^3$ one can enhance ΔT by reducing the temperature but one is also limited by V not

being too large. An optimum set of parameters is to have T in the mK range and a detector of mass of $\sim 1\text{kg}$. Detectors of this type are currently being built.

3.22 Superconducting Detectors

Here superconducting granules are suspended in a dielectric carrier in the presence of a magnetic field (Fig. 3.) The superconductors are put into a metastable

Fig. 3 Schematic diagram of a superconducting detector¹.

state and hence the magnetic field is excluded from the granules by the Meissner effect. When a \tilde{Z}_1 strikes a granule, the deposited energy triggers the transition to the normal state, and the magnetic flux movement then produces a signal in the pickup coil. The characteristic size of such detectors are also about 1 kg.

Background for these DM detectors include cosmic ray muons and natural radioactivity. The present sensitivity expected is 0.1 events/kg da, and this might be improved at a later date (i.e. by going underground) to 0.01 events/kg da.

4. Dynamical Model

To calculate the event rates expected at terrestrial detectors, we need to calculate two items: (1) the relic density of the \tilde{Z}_1 , in order to make sure that $\Omega_{\tilde{Z}_1} h^2$ lies in the range of Eq. (4), and (2) the \tilde{Z}_1 -nucleus cross section to obtain the expected event rate for a given detector. The bounds on $\Omega_{\tilde{Z}_1} h^2$ significantly limits the SUSY parameter space and hence strongly affects the event rates obtained from the \tilde{Z}_1 -nucleus cross sections.

In order to carry out the above calculations one needs a dynamical model and we use here models based on supergravity grand unification⁴. These models have

the advantages of (i) being consistent with the LEP data on unification of couplings at $M_G \simeq 10^{16}$ GeV, (ii) generating spontaneous breaking of supersymmetry at M_G (in the “hidden” sector), (iii) generating naturally spontaneous breaking of $SU(2) \times U(1)$ at the electroweak scale by radiative corrections, and (iv) having all new SUSY phenomena described by only 4 new extra parameters and one sign.

The minimal supergravity model (MSGM) is characterized at the Gut scale by a superpotential

$$W = \mu_0 H_1 H_2 + W_Y + \frac{1}{M_G} W^{(4)} \quad (5)$$

where W_Y is the cubic Yukawa couplings and $W^{(4)}$ is the quartic non-renormalizable couplings (possibly leading to proton decay). In addition there is a soft supersymmetry breaking effective potential

$$V_{SB} = m_0^2 \sum_a z_a^\dagger z_a + (A_0 W_Y + B_0 \mu_0 H_1 H_2 + h.c.) \quad (6)$$

where $\{z_a\}$ are the scalar fields, and a universal gaugino mass term $L_{mass}^\lambda = -m_{1/2} \bar{\lambda}^\alpha \lambda^\alpha$ where λ^α are the gaugino fields. The scalar mass m_0 (and cubic soft breaking constant A_0) are universal provided the agent of supersymmetry breaking in the hidden sector (e.g. the super Higgs field) communicates with the physical sector in a flavor independent way. This is automatically the case for contributions arising from the effective potential (since there the only communication is gravitational) and will be true in general if the couplings of the superHiggs to the physical fields in the Kahler potential is also flavor independent. Eq. (6) guarantees a natural suppression of unwanted FCNI.

Using the renormalization group equations (RGE) one obtains at the electroweak scale the conditions for $SU(2) \times U(1)$ breaking

$$\frac{1}{2} M_Z^2 = -\mu^2 + \frac{m_{H_1}^2 - m_{H_2}^2 \tan^2 \beta}{\tan^2 \beta - 1}; \quad \sin 2\beta = -\frac{B\mu}{2\mu^2 + m_{H_1}^2 + m_{H_2}^2} \quad (7)$$

where $\tan \beta \equiv \langle H_2 \rangle / \langle H_1 \rangle$ and $m_{H_i}^2$ are the Higgs running masses including 1-loop corrections. One can then eliminate B_0 and μ_0^2 leaving four parameters m_0 , $m_{1/2}$, A_t (the t-quark A parameter at the electroweak scale), $\tan \beta$ and the sign of μ_0 to determine all the masses and interactions of the 32 new SUSY particles. One therefore expects a large number of relations holding between the SUSY masses. If one limits the parameter space so that (i) experimental mass bounds of LEP and the

Tevatron are obeyed, (ii) $m_0, m_{\tilde{g}} < 1$ TeV, $m_{\tilde{g}}$ the gluino mass (so that no extreme fine tuning of parameters occurs), and (iii) radiative breaking of $SU(2) \times U(1)$ occurs at the electroweak scale, then the following “scaling” relations hold throughout most of the parameter space⁵: $2m_{\tilde{Z}_1} \cong m_{\tilde{Z}_2} \cong m_{\tilde{W}_1} \simeq (\frac{1}{3} - \frac{1}{4})m_{\tilde{g}}$; $m_{\tilde{Z}_3} \cong m_{\tilde{Z}_4} \cong m_{\tilde{W}_2} \gg m_{\tilde{Z}_1}$; and $m_{H^0} \cong m_{A^0} \cong m_{H^\pm} \gg m_h$. Here the \tilde{W}_i are the charginos, the \tilde{Z}_i are the neutralinos, h^0 and H^0 are the CP even Higgs, A^0 is the CP odd Higgs and H^\pm is the charged Higgs. In addition one always has $\tan\beta > 1$.

5. Calculation of \tilde{Z}_1 Relic Density

As discussed above, Eq. (4) puts a significant constraint on the allowed region of parameter space. Since this band is relatively narrow, it is important to include major effects in calculating the relic density.

R parity makes the \tilde{Z}_1 produced in the early universe absolutely stable. How-

Fig. 4 Annihilation diagrams of the \tilde{Z}_1 in the early universe.

ever, they can annihilate in pairs and the main annihilation diagrams are shown in Fig. 4. The calculation of the mass density of the \tilde{Z}_1 remaining at the present time proceeds as follows⁶: Initially the \tilde{Z}_1 are in thermal equilibrium with the background and the reactions of Fig. 4 go forward and backward. However, when the annihilation rate falls below the expansion rate of the universe, “freezeout” occurs at temperature T_f , and the \tilde{Z}_1 disconnect from the background. The \tilde{Z}_1 then continue to annihilate and the amount left at present time is⁶

$$\Omega_{\tilde{Z}_1} h^2 \cong 2.4 \times 10^{-11} \left(\frac{T_{\tilde{Z}_1}}{T_\gamma} \right)^3 \left(\frac{T_\gamma}{2.73} \right)^3 \frac{N_f}{J(x_f)} \quad (8)$$

where N_f is the effective number of degrees of freedom,

$$J(x_f) = \int_0^{x_f} \langle \sigma v \rangle dx; \quad x = kT/m_{\tilde{Z}_1}, \quad (9)$$

and $\langle \sigma v \rangle$ is the thermal average (σ = annihilation cross section, v = relative velocity):

$$\langle \sigma v \rangle = \int_0^\infty dv v^2 (\sigma v) \text{Exp}(-v^2/4x) \int_0^\infty dv v^2 \text{Exp}(-v^2/4x) \quad (10)$$

In general, the annihilation occurs non-relativistically, i.e. $x_f \approx \frac{1}{20}$. However, this does not mean one can always make a non-relativistic expansion, $\sigma v = a + bv^2 + \dots$, in performing the calculation of $J(x_f)$. As has been pointed out⁷, such an expansion fails near an s-channel pole. For example for the h pole one has

$$\sigma v = A \frac{(s - 4m_{\tilde{Z}_1}^2)/m_{\tilde{Z}_1}^2}{(s - m_h^2)^2 + m_h^2 \Gamma_h^2} \quad (11)$$

where $\Gamma_h = \mathcal{O}(\text{MeV})$ is the h boson width and A is a constant. Since Γ_h is small one must treat the pole more carefully. The danger of not doing so is shown in Fig. 5⁸.

Fig. 5. Ω_{approx} is the \tilde{Z}_1 relic density calculated in the $\sigma v = a + bv^2$ approximation and Ω is the relic density calculated rigorously using Eq. (11). The h and Z poles occur at the points where the curve descends through zero.

Note the long tail where an error of a factor of ≈ 2 can be made well past the h and Z poles. One finds, for example, that Ω_{approx} has an error of $> 25\%$ for 65% of the parameter points where $m_{\tilde{g}} < 450$ GeV, while Ω_{approx} is a good approximation for $\simeq 100\%$ of the points for $m_{\tilde{g}} \geq 450$ GeV. The reason for this is the scaling relations discussed at the end of Sec. 4. Since $2m_{\tilde{Z}_1} > (\frac{1}{3} - \frac{1}{4})m_{\tilde{g}}$, when $m_{\tilde{g}} \geq 450$ GeV, the \tilde{Z}_1 has passed both the h and Z poles (where $2m_{\tilde{Z}_1} \simeq m_h$ or M_Z). (Recall that in

the MSGM one has $m_h < 130$ GeV.) However, for lighter $m_{\tilde{g}}$ the pole effects are very important, since one is almost always near an h or a Z pole.

6. Detector Rates

The dark matter detectors discussed in Sec. 3.2 detect the \tilde{Z}_1 from their scattering by quarks in the nucleus. The basic diagrams are given in Fig. 6. They

Fig. 6 \tilde{Z}_1 -quark scattering diagrams. \tilde{q} is the squark.

are mainly crossed diagrams to the relic annihilation diagrams of Fig. 4. Thus to a rough approximation, when the relic density of the \tilde{Z}_1 is small (i.e. there is a large annihilation cross section) the number of scattering events will be large. This makes the results somewhat sensitive to the lower bound, $\Omega_{\tilde{Z}_1} h^2 \geq 0.1$, that we have chosen, and we will discuss this sensitivity below in Sec. 8.

Calculations show that it is possible to represent the \tilde{Z}_1 -q scattering by an effective Lagrangian⁹:

$$L_{eff} = (\bar{\chi}_1 \gamma^\mu \gamma^5 \chi_1) \bar{q} \gamma^\mu (A_q P_L + B_q P_R) q + (\bar{\chi}_1 \chi_1) C_q m_q \bar{q} q \quad (12)$$

where $\chi_1(x)$ is the \tilde{Z}_1 field, $q(x)$ the quark field, and m_q its mass and $P_{R,L} = \frac{1}{2}(1 \pm \gamma^5)$. The coefficients A_q , B_q arise from the Z t-channel and \tilde{q} s-channel poles while C_q arises from the h and H t-channel and \tilde{q} s-channel poles.

The first term of Eq. (12) leads to spin dependent incoherent scattering by the quarks in the nucleus, while the second term leads to coherent scattering, where the masses of all the quarks (and hence nucleons) approximately add coherently. Thus the coherent amplitude will contain a factor M_N larger than the incoherent amplitude, where M_N is the nucleus mass.

In general, the \tilde{Z}_1 field is a mix of higgsinos and gauginos:

$$\chi_1 = n_1 \tilde{W}_3 + n_2 \tilde{B} + n_3 \tilde{H}_2 + n_4 \tilde{H}_1 \quad (13)$$

where \tilde{W}_3 and \tilde{B} are the W_3 and U(1) [Bino] gauginos, and $\tilde{H}_{1,2}$ are Higgsinos. The n_i can be computed in terms of the basic SUSY parameters m_0 , $m_{\frac{1}{2}}$, A_t and $\tan\beta$. Over most of the allowed part of the parameter space one finds

$$n_2 > n_1, n_4 \gg n_3 \quad (14)$$

In SUSY theory, m_h is small i.e. $m_h < (120 - 130)$ GeV, and $m_{H^2} \gg m_h^2$. In spite of this it is remarkable that the H contribution is important in the coherent amplitude¹⁰. The reason for this is the following. One finds for C_q the result

$$C_q = \frac{g_2^2}{4M_W^2} \left[\left\{ \begin{array}{cc} \frac{\cos\alpha}{\sin\beta} & \frac{F_h}{m_h^2} \\ -\frac{\sin\alpha}{\cos\beta} & \frac{F_h}{m_h^2} \end{array} \right\} + \left\{ \begin{array}{cc} \frac{\sin\alpha}{\sin\beta} & \frac{F_H}{m_H^2} \\ \frac{\cos\alpha}{\cos\beta} & \frac{F_H}{m_H^2} \end{array} \right\} \right] \begin{array}{l} u - quark \\ d - quark \end{array} \quad (15)$$

where α is the rotation angle needed to diagonalize the 2×2 h-H mass matrix, $F_h = (n_1 - n_2 \tan\theta_W)(n_3 \cos\alpha - n_4 \sin\alpha)$ and $F_H = (n_1 - n_2 \tan\theta_W)(n_3 \sin\alpha - n_4 \cos\alpha)$. In calculating the h-H mass matrix, one must include as is well-know, the one loop corrections due to the fact that m_t is large. One finds then, for the allowed part of the parameter space that α is generally small i.e. $\alpha \approx \frac{1}{10}$. This result combined with Eq. (14) shows that generally $\cos\alpha F_H \gg \sin\alpha F_h$, which can overcome the reduction in the d-quark amplitude due to the largeness of m_H . One finds, in general, as one varies over the parameter space that the H contribution to C_q can vary from $\frac{1}{10}$ to 10 times the h contribution for d-quarks, but is generally a small correction for u-quarks.

7. Detector Event Rates

The total event rate expected for a given dark matter detector can be written in the following form⁹

$$R = [R_{coh} + R_{inc}] \left[\frac{\rho_{\tilde{Z}_1}}{0.3 \text{ GeV cm}^{-3}} \right] \left[\frac{\langle v_{\tilde{Z}_1} \rangle}{320 \text{ km/s}} \right] \frac{\text{events}}{\text{kg da}} \quad (16)$$

where

$$R_{coh} = \frac{4m_{\tilde{Z}_1} M_N^3 M_Z^4}{[M_N + m_{\tilde{Z}_1}]^2} |A_{coh}|^2 \quad (17)$$

$$R_{inc} = \frac{4m_{\tilde{Z}_1}M_N}{[M_N + m_{\tilde{Z}_1}]^2} \lambda^2 J(J+1) |A_{inc}|^2 \quad (17a)$$

Here $A_{coh} \sim C_q$, $A_{inc} \sim B_q - A_q$, J is the nuclear spin and $\lambda < N|\vec{J}|N > = < N|\sum \vec{S}_i|N >$ where the sum is over the spins of all nucleons in nucleus N . Note that for large M_N , R_{coh} increases as M_N while R_{inc} decreases as $1/M_N$. This additional M_N^2 factor in R_{coh} is as expected from the discussion following Eq. (12).

We have examined the expected detector event rates for detectors made from the following nuclei¹¹: ^3He , ^{40}Ca , ^{19}F , ^{76}Ge + ^{73}Ge (equal mix of isotopes), ^{73}Ga , ^{75}As , ^{23}Na , ^{127}I , and ^{207}Pb . Both ^{19}F and ^3He have strong spin interactions, while Ge , I and Pb are increasingly heavy and hence have increasingly strong coherent scattering.

The parameter space studied was

$$100\text{GeV} \leq m_0, m_{\tilde{g}} \leq 1\text{TeV}; -6 \leq A_t/m_0 \leq 6; 2 \leq \tan\beta \leq 20 \quad (18)$$

and the mesh used was $\Delta m_0 = 100$ GeV, $\Delta m_{\tilde{g}} = 25$ GeV, $\Delta A_t/m_0 = 0.5$, $\Delta(\tan\beta) = 2, 4$, and the t quark mass was set at $m_t = 167$ GeV.

The dependence of the event rates on the SUSY parameters is generally quite complicated. However, it is possible to understand these dependences in a qualitative way. Fig. 7 shows that event rates decrease rapidly with $m_{\tilde{g}}$. This follows from the fact that the \tilde{Z}_1 becomes increasingly Bino as $m_{\tilde{g}}$ increases i.e. n_2 of Eq. (13) grows and hence n_3 and n_4 shrinks, making the interference between the gaugino and Higgsino in F_h and F_H (needed for coherent scattering) become small. CaF_2 has the strongest spin dependent forces while Pb is the heaviest of the detectors chosen. One sees from this that R_{coh} is significantly larger than R_{inc} , a general feature.

Fig. 7 also shows that R increases with $\tan\beta$, a feature which can be seen in detail in Fig. 8. This behavior follows from the $1/\cos\beta \sim \tan\beta$ factor for d -quark scattering in Eq. (15). One sees again the scaling of R with M_N i.e. the $^{23}\text{Na}^{127}\text{I}$ curves lie above the $^{73}\text{Ge} + ^{76}\text{Ge}$ curves. The dependence of R on m_0 is somewhat complicated. One expects R to decrease with m_0 since $m_{\tilde{q}}^2$ increases with m_0^2 and hence the squark s -channel contribution of Fig. 6. shrinks. In addition, however, the μ^2 determined by the radiative breaking condition Eq. (7) increases as m_0^2 increases, making the \tilde{Z}_1 increasingly more Bino like which further reduces R_{coh} .

This decrease in R with increasing m_0 shown in Fig. 9.

Fig. 7 Event rate as a function of $m_{\tilde{g}}$ for Pb (solid) and Ca F_2 (dashed) detectors. The upperline in each pair is for $\tan\beta=20$ and lower line for $\tan\beta=6$. The curves are for $A_t/m_0=1.5$, $m_0=100$ GeV, $\mu > 0$.

Fig. 8 Event rates for NaI and Ge detectors vs. $\tan\beta$ for $m_{\tilde{g}}=275$ GeV. The dot dash curve is for $A_t/m_0=1.0$, $m_0=200$ GeV, the dashed curve for $A_t/m_0=0.5$, $m_0=300$ GeV, and the solid curve for $A_t/m_0=0.0$, $m_0=200$ GeV, $\mu > 0$. The upper curve of each pair is for NaI, the lower for Ge. The A_t, m_0 parameters were chosen so that $\Omega_{\tilde{Z}_1} h^2$ is approximately the same in each case of a fixed $\tan\beta$.

The \tilde{Z}_1 mass is an increasing function of $m_{\tilde{g}}$ and as $m_{\tilde{g}}$ increases, the relic density

Fig. 9. Event rate as a function of m_0 for $m_{\tilde{g}}=300$ GeV, $A_t/m_0=0.5$, $\tan\beta=8$, $\mu > 0$. The solid curves from bottom to top are for Ge, NaI and Pb and the dashed curve is for CaF_2 .

increases (i.e. the annihilation cross section in the early universe decreases). The upper bound on $\Omega_{\tilde{Z}_1} h^2$ then leads to an upper bound on $m_{\tilde{g}}$. In general this upper bound on $m_{\tilde{g}}$ of about 750 GeV. This effect is shown in Fig. 10.

Fig. 10 $(m_{\tilde{g}})_{max}$ vs. $\tan\beta$, $\mu > 0$. The curves, from bottom to top are for $A_t/m_0 = -0.5, 0.5, 1.0, 2.0, 3.0$ and 3.5 .

One may scan the entire parameter space to obtain the maximum and minimum event rates as a function of A_t . These are shown in Fig. 11 for $\mu > 0$ and Fig. 12

for $\mu < 0$ for the domain $2 \leq \tan\beta \leq 20$. As expected from Fig. 8, the maximum event rates occur for $\tan\beta=20$. However, the minimum rates occur at different $\tan\beta$ for different A_t . Current dark matter detectors can achieve a sensitivity of $R \simeq 0.1$

Fig. 11 Maximum and minimum event rates as a function of A_t for $\mu < 0$ for Pb (solid) and CaF_2 (dashed) detectors.

Fig. 12 Same as Fig. 11 for $\mu > 0$.

event/kg da with perhaps a factor of 10 improvement in future sensitivity. One sees that only part of the parameter space i.e. the region with relatively large $\tan\beta$ will be accessible. The detectors with large R_{coh} (e.g. Pb) are generally considerably more sensitive than those with large R_{inc} (e.g. CaF_2).

8. Sensitivity To Bounds on $\Omega_{\tilde{Z}_1} h^2$

In the preceeding discussion, we have assumed the bounds of Eq. (4) for $\Omega_{\tilde{Z}_1} h^2$. The results have some sensitivity to the choice of endpoints and we examine this here.

As discussed in Sec. 6, small $\Omega_{\tilde{Z}_1} h^2$ generally leads to large event rates, and R generally rises rapidly as $\Omega_{\tilde{Z}_1} h^2$ decreases near its lower bound. The largest R, however, occurs for small $m_{\tilde{g}}$ (see e.g. Fig. 7), and hence by the scaling relations discussed at the end of Sec. 4, for small $m_{\tilde{W}_1}$. However, there also exist cuts, e.g. $m_{\tilde{W}_1} > 45$ GeV, required by LEP data which forbids $m_{\tilde{W}_1}$ from getting too small. Thus one has a sharply rising function R hitting an experimental constraint on the parameter space. This is the origin of the sharp peaks in the maximum event rate curves of Figs. 11 and 12, i.e. the maximum event rate gets quite large or not depending on whether or not the parameter point passes the experimental cut.

The sensitivity of this effect is seen in Fig. 13 (for Pb) and Fig. 14 (for CaF₂) for the maximum event rates when one increases the minimum value of $\Omega_{\tilde{Z}_1} h^2$

Fig. 13 Maximum event rate as a function of A_t/m_0 for $\mu < 0$ for Pb detector. Solid curve is for $\Omega_{\tilde{Z}_1} h^2 > 0.10$ and dot-dash for $\Omega_{\tilde{Z}_1} h^2 > 0.15$.

from 0.10 to 0.15. The sharp peaks get clipped off when the lower bound on $(\Omega_{\tilde{Z}_1} h^2)$ is increased. There still remains, however, a sizable portion of the parameter space

where R is large.

Fig. 14 Maximum event rate as a function of A_t/m_0 for $\mu < 0$ for CaF_2 detector. Dashed curve is for $\Omega_{\tilde{Z}_1} h^2 > 0.10$ and solid curve for $\Omega_{\tilde{Z}_1} h^2 > 0.15$.

The minimum event rates are sensitive to the upper bound chosen for $\Omega_{\tilde{Z}_1} h^2$. This can be seen in Fig. 15. The minimum event rates increases by more than a factor of 10 as $(\Omega_{\tilde{Z}_1} h^2)_{max}$ is reduced from 0.35 to 0.20. One notes that the inflationary scenario favors a small Hubble constant, i.e. $h=0.5$, so that the pre-

Fig. 15 Minimum event rates as a function of $(\Omega_{\tilde{Z}_1} h^2)_{max}$ for Pb detector (solid curve) and CaF_2 detector (dashed curve), $\mu < 0$.

dicted age of the universe not be inconsistent with estimated ages of globular star

clusters. Even if one assumed $\Omega_{\tilde{Z}_1} = 1$ (i.e. no hot dark matter and negligible baryonic dark matter) this would imply $\Omega_{\tilde{Z}_1} h^2 = 0.25$. Thus the larger minimum event rates of Fig. 15 may possibly be the correct choice.

The reason the minimum event rate increases with decreasing $(\Omega_{\tilde{Z}_1} h^2)_{max}$ is that the minimum rates occur for the maximum values of $m_{\tilde{g}}$. (As seen in Fig. 7 the event rate drops with increasing $m_{\tilde{g}}$.) Further, as discussed in Sec. 7, Fig. 10, $m_{\tilde{g}}$ possesses a maximum value because of the upper bound on $\Omega_{\tilde{Z}_1} h^2$. Fig. 16 shows the dependence of $(m_{\tilde{g}})_{max}$ on the maximum value of $\Omega_{\tilde{Z}_1} h^2$. As

Fig. 16 Maximum value of $m_{\tilde{g}}$ as a function of the upper bound on $\Omega_{\tilde{Z}_1} h^2$ for $\mu < 0$, $A_t/m_0 = 0.5$, $\tan\beta = 6$. Results are insensitive to the values of A_t and $\tan\beta$.

$(\Omega h^2)_{max}$ is reduced from 0.35 to 0.20, $(m_{\tilde{g}})_{max}$ drops from 750 GeV to 400 GeV. The inflationary scenario thus favors smaller values of $m_{\tilde{g}}$. Such low $m_{\tilde{g}}$ implies that the gluino could be detected at suggested energy upgrades of the Tevatron (e.g. at the DiTevatron¹²).

9. The $b \rightarrow s + \gamma$ Decay

Recently, the CLEO Collaboration have measured the branching ratio for the inclusive decay $B \rightarrow X_s + \gamma$:

$$BR(B \rightarrow X_s + \gamma) = (2.32 \pm 0.51 \pm 0.29 \pm 0.32) \times 10^{-4} \quad (19)$$

where the first error is statistical and the last two errors systematic. Combining all errors in Gaussian quadrature, one has in the spectator approximation that

$\text{BR}(b \rightarrow s + \gamma) \cong (2.32 \pm 0.67) \times 10^{-4}$. The $b \rightarrow s + \gamma$ decay is of particular interest in that it begins at the loop level as it is a FCNC process. This means that SM and new physics effects enter at the same loop level and one could expect large [i.e. $\mathcal{O}(1)$] new physics corrections to the SM predictions. Thus the $b \rightarrow s + \gamma$ decay is an excellent process for detecting new physics. We will investigate what effects the current experimental value has on the SUSY parameter space and hence on the expected dark matter detection rates.

The elementary diagrams at the electroweak scale $\mu = M_W$ are shown in Fig. 17. The $W^- - t$ intermediate state is the Standard Model contribution, while the

Fig. 17 Elementary penguin diagrams for $b \rightarrow s + \gamma$ decay at the electroweak scale $\mu = M_W$. Only the third generation quarks and squarks make a significant contribution.

$H^- - t$ and $\tilde{W}^- - \tilde{t}$ represent additional SUSY contributions. The interactions can be represented by an effective Lagrangian for a transition magnetic dipole interaction¹³:

$$L_{eff} = G_F V_{tb} V_{ts}^* m_b A_\gamma \bar{s}_L \sigma^{\mu\nu} b_R F_{\mu\nu}, \quad (20)$$

where the coefficient A_γ can be evaluated in terms of the basic parameters $m_0, m_{\tilde{g}}, A_t, \tan\beta$. In order to calculate the decay rate, however, one must use the RG equations to evaluate the amplitude at the b scale $\mu \approx m_b$. This causes operator mixing with the gluonic transition magnetic moment operator $\bar{s}_L \sigma^{\mu\nu} T^a b_R G_{\mu\nu}^a$ (where $G_{\mu\nu}^a$, $a=1\dots 8$ is the gluon field strength) and six four quark operators. It is convenient to consider the ratio

$$R = \frac{\text{BR}(b \rightarrow s + \gamma)}{\text{BR}(b \rightarrow c + e + \bar{\nu}_e)} \cong \frac{\text{BR}(B \rightarrow X_s + \gamma)}{\text{BR}(B \rightarrow X_c + e + \bar{\nu}_e)} \quad (21)$$

since poorly known CKM matrix elements and m_b factors cancel out in the ratio. One can recover the $b \rightarrow s + \gamma$ rate then from the experimental number of the charm semi-leptonic rate: $\text{BR}(B \rightarrow X_c + e + \bar{\nu}_e) = (10.7 \pm 0.5)\%$. To leading order (LO) QCD the value of R is^{13,14}

$$R = \frac{6\alpha}{\pi} \frac{|V_{ts}^* V_{tb}|^2}{|V_{cb}|^2} \frac{\left[\eta^{\frac{16}{23}} C_7(M_W) - \frac{8}{3} (\eta^{\frac{14}{23}} - \eta^{\frac{16}{23}}) C_8(M_W) + C_2(M_W) \right]^2}{I(z) \left[1 - \frac{2}{3\pi} \frac{\alpha_3(M_Z)}{\eta} f(z) \right]} \quad (22)$$

where $\eta = \alpha_3(M_Z)/\alpha_3(m_b) = 0.548$, $z = m_c/m_b = 0.316 \pm 0.013$, $I(z)$ is a phase space factor for the $b \rightarrow c e \bar{\nu}_e$ decay, $C_7(C_8)$ are the Wilson coefficients for the photonic (gluonic) magnetic penguin operators and C_2 comes from the operator mixing with the 4-quark operators.

There are a number of theoretical uncertainties in the above calculation which can be summarized as follows: (i) Errors in input parameters i.e. $\alpha_3(M_Z)$, m_b/m_c , CKM factor, $\text{BR}(B \rightarrow X_c e \bar{\nu}_e)$; (ii) Errors in the spectator approximation; (iii) There are large NLO (next to leading order) QCD corrections. This can be estimated by letting μ vary from $m_b/2$ to $2 m_b$ and are seen to be about $\pm 25\%$; (iv) Heavy mass threshold corrections in running the RGE from the t quark/squark, H^- , \tilde{W} threshold^{15,16} down to m_b . These are about 15% for the t -quark and estimated to be $\pm 15\%$ for the SUSY thresholds¹⁶. Thus current theory has an overall error of about $\pm 30\%$.

The current CLEO measurement of the $b \rightarrow s + \gamma$ rate has a significant effect on the expected dark matter detector counting rates. Fig. 18 shows the expected $b \rightarrow s + \gamma$ rate for a characteristic choice of parameters¹⁷. The $\text{BR}(b \rightarrow s + \gamma)$ is increased when μ and A_t have the same sign relative to the value when μ and A_t have the opposite sign. (This effect comes from the $\tilde{W} - \tilde{t}$ diagram of Fig. 17). Thus regions where μ and A_t have the same sign can exceed the CLEO measured value and such regions of parameter space are then experimentally eliminated. One sees from Fig. 18 also that the $\text{BR}(b \rightarrow s + \gamma)$ is largest when m_0 and $m_{\tilde{g}}$ are small, which we also saw is the region when the dark matter counting rate R is largest. Thus one expects that the maximum values of R get eliminated when μ and A_t

have the same sign. This may be seen in Fig. 19¹⁸ (for $\mu < 0$) and Fig. 20¹⁸ (for

Fig. 18 $\text{BR}(b \rightarrow s + \gamma)$ as a function of $m_{\tilde{W}_1}$ for $\tan\beta=5.0$, $|A_t/m_0|=0.5$, $m_t=165$ GeV, $\alpha_G^{-1}=24.11$. Graphs (a) and (b) are for $A_t < 0$, (c) and (d) for $A_t > 0$, while (a) and (c) are for $\mu > 0$ and (b) and (d) for $\mu < 0$.

$\mu > 0$). There we have plotted R_{Max} for the Pb detector without the $b \rightarrow s + \gamma$ condition and R_{Max} with parameter points excluded when the predicted $b \rightarrow s + \gamma$

Fig. 19 R_{Max} vs A_t/m_0 for Pb detectors for $\mu < 0$. The solid line is the expected rate without $b \rightarrow s + \gamma$ constraint and the dashed line is the rate with parameter points excluded where the predicted $b \rightarrow s + \gamma$ rate lies outside 95% C.L. bound of the experimental value of Eq. (19).

rate exceeds the 95% C.L. of Eq. (19). We see that when μ and A_t have the same sign, the maximum event rate drops sharply, well below what could be observable

Fig. 20 Same as Fig. 19 for $\mu > 0$.

in the foreseeable future. However, the COBE bounds of Eq. (4) turn out to allow mostly $A_t > 0$. Thus the major effect of the CLEO measurement of $b \rightarrow s + \gamma$ is for $\mu > 0$ where for most of the parameter space, the event rate will be very small and hence unobservable. However, for $\mu < 0$, the $b \rightarrow s + \gamma$ measurement does not effect the expected rates very much, and large dark matter event rates are still possible.

10. Proton Decay

The preceeding discussion has been for a generic supergravity Gut model described in Sec. 3 for the parameter domain of Eq. (18). Results are generally independent of the Gut group and Gut physics provided that Gut threshold effects are not so strong that they prevent grand unification from occuring at $M_G \simeq 10^{16}$ GeV.

Proton decay is characteristic of all supergravity Gut models except for the flipped SU(5) model¹⁹. Further supersymmetry generally implies a unique dominant decay mode:

$$p \rightarrow \bar{\nu} + K^+ \quad (23)$$

Thus the observation of this decay would not only indicate the validity of grand unification but also of supersymmetry. One can suppress this decay by specially tailoring the form of the Gut Higgs sector, but this generally requires some awkward fine tuning. We consider here “SU(5)-type” proton decay²⁰ which arises via the

exchange of a superheavy color triplet Higgsino \tilde{H}_3 , Fig. 21. (This can happen in SU(5), O(10), E_6 etc. Gut groups.) The decay rate can be written as,

$$\Gamma(p \rightarrow \bar{\nu} + K^+) = \frac{const}{M_{H_3}^2} |B|^2, \quad (24)$$

where B is the loop amplitude. The current experimental bound²¹,

Fig. 21 Example of proton decay diagram for $p \rightarrow \bar{\nu} + K^+$ for supersymmetric grand unification. The \tilde{H}_3 vertices violate baryon and lepton number.

$\tau(p \rightarrow \bar{\nu} + K) > 1.0 \times 10^{32}$ y, then leads to a bound on B

$$B \lesssim 100 (M_{H_3}/M_G) GeV^{-1}; \quad M_G = 2 \times 10^{32} GeV \quad (25)$$

We restrict M_{H_3} to obey $M_{H_3}/M_G < 10$ in the following so that the Gut scale be disjoint from the Planck scale. (For larger M_{H_3} one might expect large Planck physics corrections to enter, the nature of which are not known.) B can be expressed in terms of $m_{\tilde{W}}, m_{\tilde{q}}$ etc. and hence by the RG equations in terms of the four basic parameters $m_0, m_{\tilde{q}}, A_t, \tan\beta$ and the sign of μ . Thus the condition on B is a constraint on the parameter space.

The second generation dominates the loop of Fig. 21 and to a rough approximation on has

$$B_2 \approx -\frac{2\alpha_2}{\alpha_3 \sin 2\beta} \frac{m_{\tilde{g}}}{m_{\tilde{q}}^2} \times 10^6 GeV^{-1} \quad (26)$$

where $m_{\tilde{q}}^2 \cong m_0^2 + 0.6m_{\tilde{g}}^2$. Thus the proton decay bounds of Eq. (25) imply small $m_{\tilde{g}}$, large m_0 e.g. $m_0 > m_{\tilde{g}}$, and small $\tan\beta$ i.e. $\tan\beta \leq 10$. One may satisfy both the COBE constraints on $\Omega_{\tilde{Z}_1} h^2$ of Eq. (4) simultaneously with the proton decay

constraints of Eq. (25) even though m_0 must be large (which usually is a region of small relic \tilde{Z}_1 annihilation). This is possible since, as discussed in Sec. 5, when $m_{\tilde{g}} \leq 450$ GeV, (the region also required by Eq. (25) for proton decay) large relic \tilde{Z}_1 annihilation can occur due to the presence of h or Z poles. In fact, one finds that in the vicinity of these poles, when m_0 is small (e.g. $m_0 < m_{\tilde{g}}$) too much annihilation occurs [i.e. $\Omega_{\tilde{Z}_1} h^2 = O(10^{-2})$] and one must increase m_0 to satisfy the lower bound of Eq. (4). These are then the domains that also satisfy the proton decay constraint. Thus in order to find the parameter space region which simultaneously satisfies the dark matter and proton decay constraints, it is essential to treat the calculation of \tilde{Z}_1 relic density in the accurate fashion discussed in Sec. 5. One finds then that the parameter points satisfying the simultaneous constraints require

$$\begin{aligned} m_{\tilde{g}} &\leq 375 \text{ GeV}; \quad m_0 \geq 500 \text{ GeV}; \quad \tan\beta \leq 10 \\ 0.0 &\leq A_t/m_0 \leq 0.5 \end{aligned} \tag{27}$$

for a t quark of mass 165 GeV.

One may next ask whether the CLEO measurement of the $b \rightarrow s + \gamma$ decay effects this result. One finds, however, that 95% of the parameter points which simultaneously satisfy the dark matter and proton decay constraints, predict a $b \rightarrow s + \gamma$ branching ratio in the LO that is within the 90% C.L. of the experimental value of Eq. (19). Thus, at the current experimental and theoretical accuracy, the $b \rightarrow s + \gamma$ decay does not effect the proton decay predictions.

However, the proton decay constraint does effect the expected dark matter event rates, and the maximum event rates are significantly reduced since $\tan\beta$ is small and m_0 is large. Parameter points which simultaneously satisfy COBE, the $b \rightarrow s + \gamma$ and proton decay constraints lead to event rates of size $R = O(10^{-3} - 10^{-4})$ events/kg da. Thus if the next round of proton decay experiments (Super Kamiokande, ICARUS) were to actually detect proton decay, the present theory implies that \tilde{Z}_1 dark matter is beyond the ability of current detectors to discover by direct detection.

11. Conclusions

Studying non high energy accelerator phenomena such as dark matter, proton decay, the $b \rightarrow s + \gamma$ decay is useful in limiting the SUSY parameter space. We have examined within the framework of the minimal supergravity model (MSGM) the ability of direct detection of dark matter when the relic density obeys Eq. (4).

- (i) The detectors which are most sensitive to coherent \tilde{Z}_1 scattering (e.g. the Pb detector) are better than the detectors most sensitive to incoherent (spin-dependent) scattering, and the heavier the nucleus, the more sensitive the detector.
- (ii) With a future sensitivity of $R > 0.01$ events/kg da, a reasonable amount, though not all, of the parameter space will be accessible to dark matter detectors.
- (iii) Raising the lower bound on $\Omega_{\tilde{Z}_1} h^2$ decreases the maximum event rate (Fig. 13, 14), and lowering the upper bound increases the minimum event rate (Fig. 15). The upper bound, $\Omega_{\tilde{Z}_1} h^2 = 0.35$, also determines an upper bound on $m_{\tilde{g}}$ of 750 GeV. Thus if this upper bound is lowered (as suggested by the inflationary scenario without a cosmological constant) the gluino would become more accessible to accelerator detection.
- (iv) The largest event rates occur at large $\tan\beta$ and small $m_{\tilde{g}}$ and small m_0 (Figs. 7-9). Thus detectors are most sensitive to these domains. Models with very large $\tan\beta$ (i.e. $\tan\beta \approx 50$) may therefore be testable by planned detectors.
- (v) The predicted $b \rightarrow s + \gamma$ decay is large in the same region of parameter space (small $m_{\tilde{g}}$, small m_0) where dark matter event R is large when μ and A_t have the same sign. The current experimental rate for this decay thus eliminates part of the parameter space where R is large. Since the relic density constraint eliminates most of the $A_t < 0$ part of the parameter space, the $b \rightarrow s + \gamma$ constraint significantly reduces the expected event rate for $\mu > 0$, but does not effect the $\mu < 0$ event rates a great deal, and large event rates can still occur for $\mu < 0$.
- (vi) For models possessing in addition SU(5)-type proton decay, there remain regions in parameter space, Eq. (27), satisfying both the relic density constraint and current proton decay bounds. These points also are within the 90% C.L. bounds of the current $b \rightarrow s + \gamma$ decay rate. However, since proton decay favors large m_0 and $\tan\beta \leq 10$, the predicted dark matter event rates are all for $R < 0.01$ event/kg da. Hence models with SU(5)-type proton decay predict that \tilde{Z}_1 dark matter will be inaccessible to current detectors.

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